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Controlling inelastic cotunneling through an interacting quantum dot by a circularly-polarized field

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We study inelastic cotunneling through a strong Coulomb-blockaded quantum dot subject to a static magnetic field and a perpendicular circularly-polarized magnetic field using a quantum Langevin equation approach. Our calculation predicts an interesting controllable cotunneling current characteristic, splitting-zero-anomaly-splitting transition of the differential conductance with increasing the driving frequency, ascribing to the role of *photon-assisted spin-flip* cotunneling processes.

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Coherent control of spin dynamics in a quantum dot (QD) using circularly-polarized magnetic field (CPF)^{1,2} in a solid-state environment has been a subject of increasing interest in recent years.³ This interest has led to a large amount of work on detection of the electron spin resonance (ESR) in a QD,^{4,5,6,7} and initializing the state of electron spin via optical absorption of circularly polarized light^{8,9,10}. Recently, the single-electron ESR in a QD-lead system with sizable Zeeman splitting has been theoretically reported to carry pure spin flow, which can be used as a fundamental element and/or a spin source device in the spintronic circuit.^{11,12} These papers have studied the transport through a QD in resonant tunneling regime by the nonequilibrium Green's function (NGF)¹¹ and the quantum rate equation approach,¹² respectively. The effect of the s-d exchange interaction between the localized spin and conduction electrons on electron transmission probability in linear transport regime has been exploited for a mesoscopic ring embedded with an magnetic impurity by the NGF.¹³ In this paper, we focus our studies, for the first time, on the influence of a driving CPF upon cotunneling (non-resonant coherent tunneling) through a strongly interacting QD in the weak-tunneling limit,^{14,15} predicting that the cotunneling current can be easily controlled by tuning the driving frequency.

Cotunneling through a single-level ϵ_d interacting QD subject to an ambient constant magnetic field \mathbf{B}_0 along the z -axis and a driving external (magnetic) CPF with frequency ω_c whose direction rotates in the plane perpendicular to the z -axis, $\mathbf{B}_1(t) = B_1(\cos \omega_c t, \sin \omega_c t, 0)$, can be described by the Hamiltonian $H = H_B + H_0 + H_1$:¹⁴

$$H_B = \sum_{\eta\mathbf{k}\sigma} \epsilon_{\eta\mathbf{k}} c_{\eta\mathbf{k}\sigma}^\dagger c_{\eta\mathbf{k}\sigma}, \quad (1a)$$

$$H_0 = -\Delta_0 S^z - \frac{1}{2} \Delta (e^{i\theta} S^+ + e^{-i\theta} S^-), \quad (1b)$$

$$H_1 = \sum_{\eta\eta', \mathbf{k}\mathbf{k}'} J_{\eta\eta'} [(c_{\eta\mathbf{k}\uparrow}^\dagger c_{\eta'\mathbf{k}'\uparrow} - c_{\eta\mathbf{k}\downarrow}^\dagger c_{\eta'\mathbf{k}'\downarrow}) S^z$$

$$+ c_{\eta\mathbf{k}\uparrow}^\dagger c_{\eta'\mathbf{k}'\downarrow} S^- + c_{\eta\mathbf{k}\downarrow}^\dagger c_{\eta'\mathbf{k}'\uparrow} S^+] + H_{\text{dir}}, \quad (1c)$$

$$H_{\text{dir}} = J_0 \sum_{\sigma} (c_{L\mathbf{k}\sigma}^\dagger + c_{R\mathbf{k}\sigma}^\dagger) (c_{L\mathbf{k}\sigma} + c_{R\mathbf{k}\sigma}), \quad (1d)$$

where $\theta = \omega_c t$, $c_{\eta\mathbf{k}\sigma}^\dagger$ ($c_{\eta\mathbf{k}\sigma}$) is the creation (annihilation) operator for electrons with momentum \mathbf{k} , spin- σ and energy $\epsilon_{\eta\mathbf{k}}$ in lead η ($= L, R$), $\mathbf{S} \equiv (S^x, S^y, S^z)$ are Pauli spin operators of electrons in the QD [$S^\pm \equiv S^x \pm iS^y$], and $J_{\eta\eta'}$ is the exchange coupling constant. $\Delta_0 = g\mu_B B_0$ is the static magnetic-field B_0 -induced Zeeman term, and $\Delta = g\mu_B B_1$ describes the spin-flip scattering caused by the CPF. H_{dir} is the potential scattering term, which is decoupled from the electron spin due to the number of electrons in the dot level being one. As a result, this term has no influence on the dynamical evolution of the electron spin and behaves only as a direct bridge to connect the left and right leads. For an Anderson model with symmetrical coupling to the leads t , we have $J_{\eta\eta'} = 2J_0 = t^2/\epsilon_d$. As in our previous paper¹⁶, we can rewrite the tunneling term, Eq. (1c), as a sum of three products of two variables: $H_1 = Q^z S^z + Q^+ S^- + Q^- S^+ + Q^{\hat{1}}$ with the same definitions of $Q^{z(\pm)}$ as in Ref. 16 and $Q^{\hat{1}} = H_{\text{dir}}$.

It is noted that our model describes the non-resonant tunneling through a localized magnetic impurity involving s-d exchange interaction with the conduction electrons in the presence of both a static magnetic field \mathbf{B}_0 and a rotating magnetic field \mathbf{B}_1 . We assume the exchange interaction to be weak so that no Kondo effect emerges, and assume that charge fluctuation completely vanishes.

It is well-known that the isolated spin-1/2 electron under the influence of both \mathbf{B}_0 and an external CPF $\mathbf{B}_1(t)$ (the Rabi problem¹), Eq. (1b), has an analytical solution.^{2,17} In the rotating frame, $x = \cos \theta S^x - \sin \theta S^y$, $y = \sin \theta S^x + \cos \theta S^y$, and $z = S^z$, the QD Hamiltonian Eq. (1b) takes the form

$$\tilde{H}_0 = -\Delta x - \delta z, \quad (2)$$

with $\delta = \Delta_0 - \omega_c$. Correspondingly, the Heisenberg equations of motion (EOM's) of these free rotating coordinates become: $\dot{x} = \delta y$, $\dot{y} = -\delta x + \Delta z$, and $\dot{z} = -\Delta y$.

Solving these resulting ordinary differential equations, we obtain the free evolutions of these rotating coordinates:

$$\begin{aligned} \mathbf{r}(t') &= \mathbf{M}(\tau)\mathbf{r}(t), \quad [\mathbf{r}(t) \equiv (x(t), y(t), z(t))^T], \quad (3) \\ \mathbf{M}(\tau) &= \begin{pmatrix} \frac{\delta^2}{\Omega^2}a^+ + \frac{\Delta^2}{\Omega^2} & -i\frac{\delta}{\Omega}a^- & -\frac{\delta\Delta}{\Omega^2}(a^+ - 1) \\ i\frac{\delta}{\Omega}a^- & a^+ & -i\frac{\Delta}{\Omega}a^- \\ -\frac{\delta\Delta}{\Omega^2}(a^+ - 1) & i\frac{\Delta}{\Omega}a^- & \frac{\Delta^2}{\Omega^2}a^+ + \frac{\delta^2}{\Omega^2} \end{pmatrix}, \end{aligned}$$

with $\tau = t - t'$, $a^\pm = \frac{1}{2}(e^{-i\Omega\tau} \pm e^{i\Omega\tau})$, and the Rabi frequency $\Omega = \sqrt{\Delta^2 + \delta^2}$.

Furthermore, the transformed interacting Hamiltonian, Eq. (1c), reads in terms of these rotating frame:

$$\tilde{H}_I = Q^x x + Q^y y + Q^z z + Q^{\hat{I}}, \quad (4)$$

with $Q^x = e^{-i\theta}Q^- + e^{i\theta}Q^+$ and $Q^y = i(e^{-i\theta}Q^- - e^{i\theta}Q^+)$. The Heisenberg EOM's for the spin operators are:

$$\dot{x} = \delta y - Q^z y + Q^y z, \quad (5a)$$

$$\dot{y} = -\delta x + \Delta z + Q^z x - Q^x z, \quad (5b)$$

$$\dot{z} = -\Delta y - Q^y x + Q^x y. \quad (5c)$$

It is clear that the spin dynamics, apart from free evolutions, are perturbatively modified by the weak tunnel coupling. To obtain the modified dynamics, we employ a generic quantum Langevin equation approach.^{16,18,19} In our derivation, operators of the QD spin and the reservoirs are first expressed formally by integration of their Heisenberg EOM's, Eq. (5), exactly to all orders of $J_{\eta\eta'}$. Next, under the assumption that the time scale of decay processes is much slower than that of free evolutions, we replace the time-dependent operators involved in the integrals of these EOM's approximately in terms of their free evolutions, Eq. (3). Thirdly, these EOM's are expanded in powers of $J_{\eta\eta'}$ up to second order. To this end, we can establish the Bloch-type dynamical equations for the averaged spin variables \mathbf{r} as:

$$\begin{pmatrix} \dot{x} \\ \dot{y} \\ \dot{z} \end{pmatrix} = \begin{pmatrix} -\Gamma_{xx} & \delta & \Gamma_{xz} \\ -\delta & -\Gamma_{yy} & \Delta \\ \Gamma_{zx} & -\Delta & -\Gamma_{zz} \end{pmatrix} \begin{pmatrix} x \\ y \\ z \end{pmatrix} + \begin{pmatrix} \gamma_x \\ 0 \\ \gamma_z \end{pmatrix}, \quad (6)$$

in which

$$\begin{aligned} \Gamma_{xx} &= 2\left(\frac{\Delta}{\Omega}\right)^2 C(\Omega) + 2\left(\frac{\delta}{\Omega}\right)^2 C(0) \\ &+ \left(1 - \frac{\delta}{\Omega}\right) C(\omega_c - \Omega) + \left(1 + \frac{\delta}{\Omega}\right) C(\omega_c + \Omega), \quad (7a) \end{aligned}$$

$$\Gamma_{xz} = 2\frac{\delta\Delta}{\Omega^2}[C(0) - C(\Omega)], \quad (7b)$$

$$\begin{aligned} \gamma_x &= \frac{\Delta}{2\Omega} \left[2R(\Omega) - 2\frac{\delta}{\Omega}R(\omega_c) - \left(1 - \frac{\delta}{\Omega}\right)R(\omega_c - \Omega) \right. \\ &\quad \left. + \left(1 + \frac{\delta}{\Omega}\right)R(\omega_c + \Omega) \right], \quad (7c) \end{aligned}$$

$$\Gamma_{yy} = 2\left(\frac{\Delta}{\Omega}\right)^2 [C(\Omega) + C(\omega_c)] + 2\left(\frac{\delta}{\Omega}\right)^2 C(0)$$

$$+ \frac{\delta}{\Omega} \left[\left(1 + \frac{\delta}{\Omega}\right) C(\omega_c + \Omega) - \left(1 - \frac{\delta}{\Omega}\right) C(\omega_c - \Omega) \right], \quad (7d)$$

$$\begin{aligned} \Gamma_{zx} &= 2\frac{\delta\Delta}{\Omega^2}C(\omega_c) + \frac{\Delta}{\Omega} \left[\left(1 - \frac{\delta}{\Omega}\right) C(\omega_c - \Omega) \right. \\ &\quad \left. - \left(1 + \frac{\delta}{\Omega}\right) C(\omega_c + \Omega) \right], \quad (7e) \end{aligned}$$

$$\begin{aligned} \Gamma_{zz} &= 2\left(\frac{\Delta}{\Omega}\right)^2 C(\omega_c) + \left(1 - \frac{\delta}{\Omega}\right)^2 C(\omega_c - \Omega) \\ &+ \left(1 + \frac{\delta}{\Omega}\right)^2 C(\omega_c + \Omega), \quad (7f) \end{aligned}$$

$$\begin{aligned} \gamma_z &= \frac{1}{2} \left[2\left(\frac{\Delta}{\Omega}\right)^2 R(\omega_c) + \left(1 - \frac{\delta}{\Omega}\right)^2 R(\omega_c - \Omega) \right. \\ &\quad \left. + \left(1 + \frac{\delta}{\Omega}\right)^2 R(\omega_c + \Omega) \right], \quad (7g) \end{aligned}$$

with the reservoir correlation function, $C(\omega)$, and the response function, $R(\omega)$, defined as

$$\begin{aligned} C(\omega) &= \frac{\pi}{2} (g_{LL} + g_{RR}) T \varphi\left(\frac{\omega}{T}\right) \\ &+ \frac{\pi}{2} g_{LR} T \left[\varphi\left(\frac{\omega + V}{T}\right) + \varphi\left(\frac{\omega - V}{T}\right) \right], \quad (8a) \end{aligned}$$

$$R(\omega) = \frac{\pi}{2} (g_{LL} + g_{RR} + 2g_{LR}) \omega, \quad (8b)$$

with $g_{\eta\eta'} \equiv J_{\eta\eta'}^2 \rho_0^2$ (ρ_0 is the constant density of states of both electrodes) and $\varphi(x) \equiv x \coth(x/2)$. V is the bias-voltage applied symmetrically to the two electrodes, $\mu_L = -\mu_R = V/2$. We use units with $\hbar = k_B = e = 1$.

Notice that these quantities of $\Gamma_{\alpha\beta}$ describe the dissipation of cotunneling processes to the dynamics of the QD spin, in conjunction with the effect of CPF: transition between two spin-splitting states $\pm\Omega/2$ (spin analog of charge optical Stark effect) due to photo-absorption and/or emission. Also, it is easy to see that the Bloch equations, Eq. (6), is exactly reduced to the results of Ref. 16 in the case of vanishing CPF, $\Delta = 0$ and $\omega_c = 0$.

The nonequilibrium steady-state spin projections of the QD in the rotating frame, $\mathbf{r}^\infty = (x^\infty, y^\infty, z^\infty)$, can be readily obtained from Eq. (6). As is evident, $x^\infty = y^\infty = 0$, while $z^\infty = \frac{1}{2}R(\Delta_0)/C(\Delta_0)$ is identical to the previous theoretical result of nonequilibrium magnetization^{16,20} in absence of driving CPF. On the other hand, in the case of vanishing static magnetic field, $\Delta_0 = 0$, the nonzero driving CPF can induce an additional spin orientation $z^\infty \neq 0$ along the rotating direction of the CPF, together with the nonzero x^∞ (not shown here). At equilibrium, this optically-induced spin orientation phenomenon has historically been termed, in literature, as either the inverse Faraday effect²¹ or as the Zeeman light shift,²² which is ascribed to the CPF-induced spin-splitting (ac spin Stark effect). In addition, one can interestingly observe a nearly vanishing y -component of the spin polarization even at CPF driving.

For nonzero static magnetic field, previous theoretical studies show that coherent suppression of tunneling may

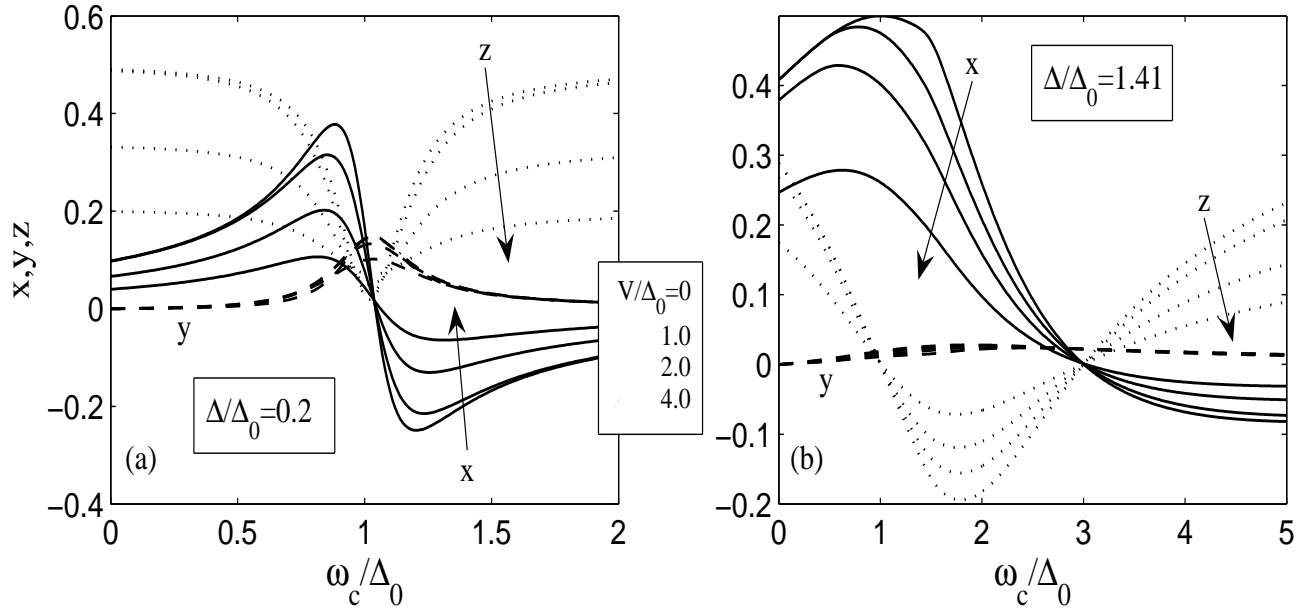


FIG. 1: The nonequilibrium spin projection \mathbf{r}^∞ vs. ω_c/Δ_0 ($\Delta_0 = 1$) with increasing bias-voltage under a weak driving CPF, $\Delta = 0.2\Delta_0$ (a), and a strong driving field, $\Delta = \sqrt{2}\Delta_0$ (b). The arrow indicates the direction of increase of bias-voltage V . The other parameters are: $g_{LL} = g_{RR} = g_{LR} = 0.05$ and $T = 0.02$.

take place only if the Rabi frequency Ω matches with the driving frequency ω_c , which corresponds to the following relationship of the driving field²⁴:

$$\Omega^* = \omega_c = \frac{\Delta_0^2 + \Delta^2}{2\Delta_0}. \quad (9)$$

In Fig. 1, we plot the nonequilibrium spin projection \mathbf{r}^∞ under the influence of both nonzero static magnetic field and driving CPFs. Different from the results of $\Delta_0 = 0$, we find an obvious nonzero polarization of the y -component of QD spin near the resonant field, $\omega_c = \Delta_0$, for the weak driving field $\Delta = 0.2\Delta_0$ [Fig. 1(a)]. Interestingly, we observe $z^\infty = 0$ at $\omega_c = \Delta_0$ for any strength of CPF, meaning that the resonant frequency of a CPF may overcome the Zeeman-splitting of a static magnetic field. At the same time, x^∞ becomes unpolarized for weak driving field [Fig. 1(a)], while reaches the maximum value for strong driving field, which decreases with increasing of the bias-voltage [Fig. 1(b)]. Moreover, we notice that (1) $x^\infty = z^\infty = 0$ at $\omega_c = 2\Omega^*$; and (2) z^∞ is negative if $\Delta_0 < \omega_c < 2\Omega^*$ for strong driving CPF.

We now proceed with the calculation of tunneling current. The current operator through the QD is defined as the time rate of change of charge density $N_\eta = \sum_{\mathbf{k},\sigma} c_{\eta\mathbf{k}\sigma}^\dagger c_{\eta\mathbf{k}\sigma}$ in lead η (we choose the left lead as an example):

$$\begin{aligned} J_L(t) &= -\dot{N}_L = i[N_L, H]_- \\ &= i(Q_{LR}^{z\uparrow\uparrow} - Q_{RL}^{z\uparrow\uparrow})S^z - i(Q_{LR}^{z\downarrow\downarrow} - Q_{RL}^{z\downarrow\downarrow})S^z \\ &\quad + i(Q_{LR}^{+} - Q_{RL}^{+})S^+ + i(Q_{LR}^{+} - Q_{RL}^{+})S^- \\ &\quad + i(Q_{LR}^{z\uparrow\uparrow} - Q_{RL}^{z\uparrow\uparrow} + Q_{LR}^{z\downarrow\downarrow} - Q_{RL}^{z\downarrow\downarrow}), \end{aligned} \quad (10)$$

where the definitions of $Q_{\eta\eta'}^{z\sigma\sigma}$ and $Q_{\eta\eta'}^\pm$ can be found in our previous paper¹⁶ and those terms in the last line are stemming from the direct tunneling term $H_{\text{dir}} = J_0 \sum_{\eta\eta',\sigma} Q_{\eta\eta'}^{z\sigma\sigma}$. The linear-response theory gives

$$I = \langle J_L(t) \rangle = -i \int_{-\infty}^t dt' \langle [J_L(t), H_I(t')]_- \rangle_0, \quad (11)$$

where the statistical average $\langle \cdots \rangle_0$ is performed with respect to decoupled two subsystems, QD and reservoirs. Inserting Eqs. (1c) and (10) into Eq. (11), one can derive the explicit expression for steady-state current in terms of the steady-state spin projections \mathbf{r}^∞ through a lengthy but straightforward calculation. In particular, we note

$$\langle [(Q_{LR}^{z\uparrow\uparrow} - Q_{RL}^{z\downarrow\downarrow})S^z, H_{\text{dir}}]_- \rangle_0 = 0, \quad (12)$$

$$\langle [(Q_{RL}^{z\uparrow\uparrow} - Q_{LR}^{z\downarrow\downarrow})S^z, H_{\text{dir}}]_- \rangle_0 = 0, \quad (13)$$

$$\langle [Q_{\eta\eta'}^\pm, S^\mp, H_{\text{dir}}]_- \rangle_0 = 0, \quad (14)$$

and

$$\begin{aligned} &\langle [(Q_{LR}^{z\uparrow\uparrow} - Q_{RL}^{z\uparrow\uparrow} + Q_{LR}^{z\downarrow\downarrow} - Q_{RL}^{z\downarrow\downarrow}), H_I] \rangle_0 \\ &= \langle [(Q_{LR}^{z\uparrow\uparrow} - Q_{RL}^{z\uparrow\uparrow} + Q_{LR}^{z\downarrow\downarrow} - Q_{RL}^{z\downarrow\downarrow}), H_{\text{dir}}] \rangle_0 \neq 0, \end{aligned} \quad (15)$$

because the QD is connected to two normal leads in the system under consideration.¹⁶ These results indicate that the contribution of the direct tunneling term to the current is independent of the dynamics of the QD and only depends on the bias-voltage and temperature of the two electrodes. To this end, the current I is²³

$$\frac{I}{\pi g_{LR}} = 4V - T \left\{ \frac{1}{2} \left[\left(\frac{\delta}{\Omega} \right)^2 + 1 \right] \mathcal{I}^+ + \frac{\delta}{\Omega} \mathcal{I}^- + \left(\frac{\Delta}{\Omega} \right)^2 \left[\varphi \left(\frac{\omega_c + V}{T} \right) - \varphi \left(\frac{\omega_c - V}{T} \right) \right] \right\} z^\infty \\ - 2T \left\{ \frac{\delta \Delta}{4\Omega^2} \mathcal{I}_c^+ + \frac{\Delta}{4\Omega} \mathcal{I}_c^- - \frac{\delta \Delta}{2\Omega^2} \left[\varphi \left(\frac{\omega_c + V}{T} \right) - \varphi \left(\frac{\omega_c - V}{T} \right) \right] + \frac{\Delta}{2\Omega} \left[\varphi \left(\frac{\Omega + V}{T} \right) - \varphi \left(\frac{\Omega - V}{T} \right) \right] \right\} x^\infty, \quad (16)$$

with

$$\mathcal{I}^\pm = \varphi \left(\frac{\omega_c + \Omega + V}{T} \right) \pm \varphi \left(\frac{\omega_c - \Omega + V}{T} \right) \\ - \left[\varphi \left(\frac{\omega_c + \Omega - V}{T} \right) \pm \varphi \left(\frac{\omega_c - \Omega - V}{T} \right) \right]. \quad (17)$$

Different from the previous results of the nondriving QD system,¹⁶ the nonzero x -component of the stationary spin polarization has additional contribution to the current besides the z -component. The linearly bias-voltage-dependent term in Eq. (16) stems from the cotunneling processes, in which the spin projection remain unchanged. The CPF has no effect on these processes. More importantly, one observes that the current formula Eq. (16) includes some new factors involving $\omega_c \pm \Omega \pm V$ [Eq. (17)], which can be intuitively ascribed to the *photo-assisted* spin-flip cotunneling processes. For instance, the term involving $\Omega - \omega_c$ results from the contribution of cotunneling accompanied by an *absorption* of one-photon [Fig. 2(a)], while the term involving $\Omega + \omega_c$ corresponds to the process assisted by a spontaneous *emission* of one-photon [Fig. 2(b)]. Therefore, we simply name these cotunnelings as $\Omega \pm \omega_c$ processes respectively below. Besides, the factors $\Omega \pm V$ in Eq. (16) stem from the spin-flip cotunneling events without spontaneous emission or absorption of photon, while the role of driving field is reflected through the effective spin-splitting Ω . In the following, we will see that the joint effects of these new terms are responsible for the controllable patterns of cotunneling through a QD by the presence of CPF.

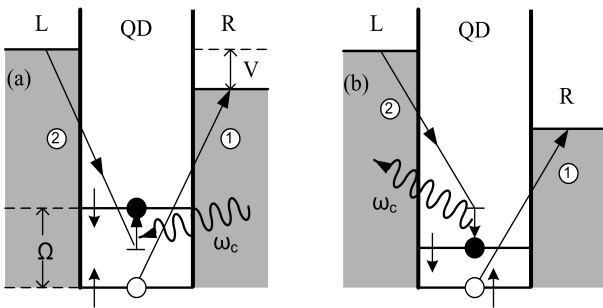


FIG. 2: Schematic diagrams of photon-absorption- ($\Omega - \omega_c$) (a) and photon-emission- ($\Omega + \omega_c$) (b) assisted spin-flip cotunneling processes.

We plot the differential conductance $G = \frac{\partial I}{\partial V}$ (in unit of $2\pi g_{LR}$) vs bias-voltage as functions of different driving frequencies ω_c for a weak driving CPF $\Delta/\Delta_0 = 0.2$

(a,b) and a strong driving field $\Delta/\Delta_0 = \sqrt{2}$ (c,d) in Fig. 3. If $\omega_c = 0$, the I - V curve reduces to the ordinary cotunneling characteristics under an ambient constant magnetic field:¹⁶ a jump at $V = \pm\Omega = \pm\sqrt{\Delta_0^2 + \Delta^2}$ [$\approx \Delta_0$ in Fig. 3(a) due to $\Delta \ll \Delta_0$ and $= \sqrt{3}\Delta_0$ in Fig. 3(c), respectively], ascribing to the energetic inactivation (activation) of spin-flip processes at bias-voltage, $|V| < (>)\Omega$. While for nonzero driving frequency, such as $\omega_c/\Delta_0 = 0.5$, the CPF effectively suppresses spin-splitting $\Omega/\Delta_0 \approx 0.54$ at the case of weak driving field, thus leading to two jumps, one of which is located at $V = \pm\Omega$ due to the pure activated spin-flip process, and another of which appears at $V = \pm(\Omega + \omega_c) \approx \pm\Delta_0$, which can be ascribed to opening of an additional channel for electron transfer cotunneling due to the one-photon-emission-assisted spin-flip resonant event. In contrast, the strong driving field $\Delta = \sqrt{2}\Delta_0$ makes $\Omega = 1.5\Delta_0$ and generates three jumps: the first one corresponds to the $\Omega - \omega_c$ process; the second one is due to the pure electronic spin-flip event (Ω); and the third one results from the $\Omega + \omega_c$ process. When the frequency increases to a bit higher than the resonant frequency [$1.0\Delta_0$ ($\Omega^* = 1.5\Delta_0$) for the weak (strong) driving field], a zero-bias peak emerges for both cases, which can be understood by the fact that the proper high frequency CPF has enough energy to spur the spin-flip cotunneling even at equilibrium [see Fig. 2(a)], while the rising voltage, $V \geq \pm\Omega$ [Figs. 3(a, b)] or $\pm(\Omega - \omega_c)$ [Figs. 3(c, d)], contrarily suppresses its activation until a pure bias-driven spin-flip process is excited at $V = \pm\Delta_0$ or $\pm\Omega$. Recently, the zero-anomaly (ZA) behavior of cotunneling current has been also reported for a QD connected to two anti-parallel ferromagnetic electrodes.^{25,26} Here, it should be pointed out that the appearance of ZA in the present system is due to photon-assisted *inelastic* spin-flip scattering becoming resonant in the presence of external magnetic field, which is different from the previous mechanism, *elastic* spin-flip event in the case of polarization leads without any magnetic field.^{25,26} As a result, the present system has more rich transport features by tuning strength and frequency of driving CPF.

For the weak CPF field with $\omega_c \approx 2\Delta_0$, the photon-assisted excitation of spin-flip cotunneling merges coincidentally into pure bias excitation, leading to the disappearance of ZA behavior [Fig. 3(b)]. While for strong CPF field [Fig. 3(d)], the ZA also vanishes if $\omega_c \geq 2\Omega^*$, which is due to the peculiar features of the nonequilibrium spin projections x^∞ and z^∞ [see Fig. 1(b)]. For higher driving frequency, the cotunneling exhibits splitting differential

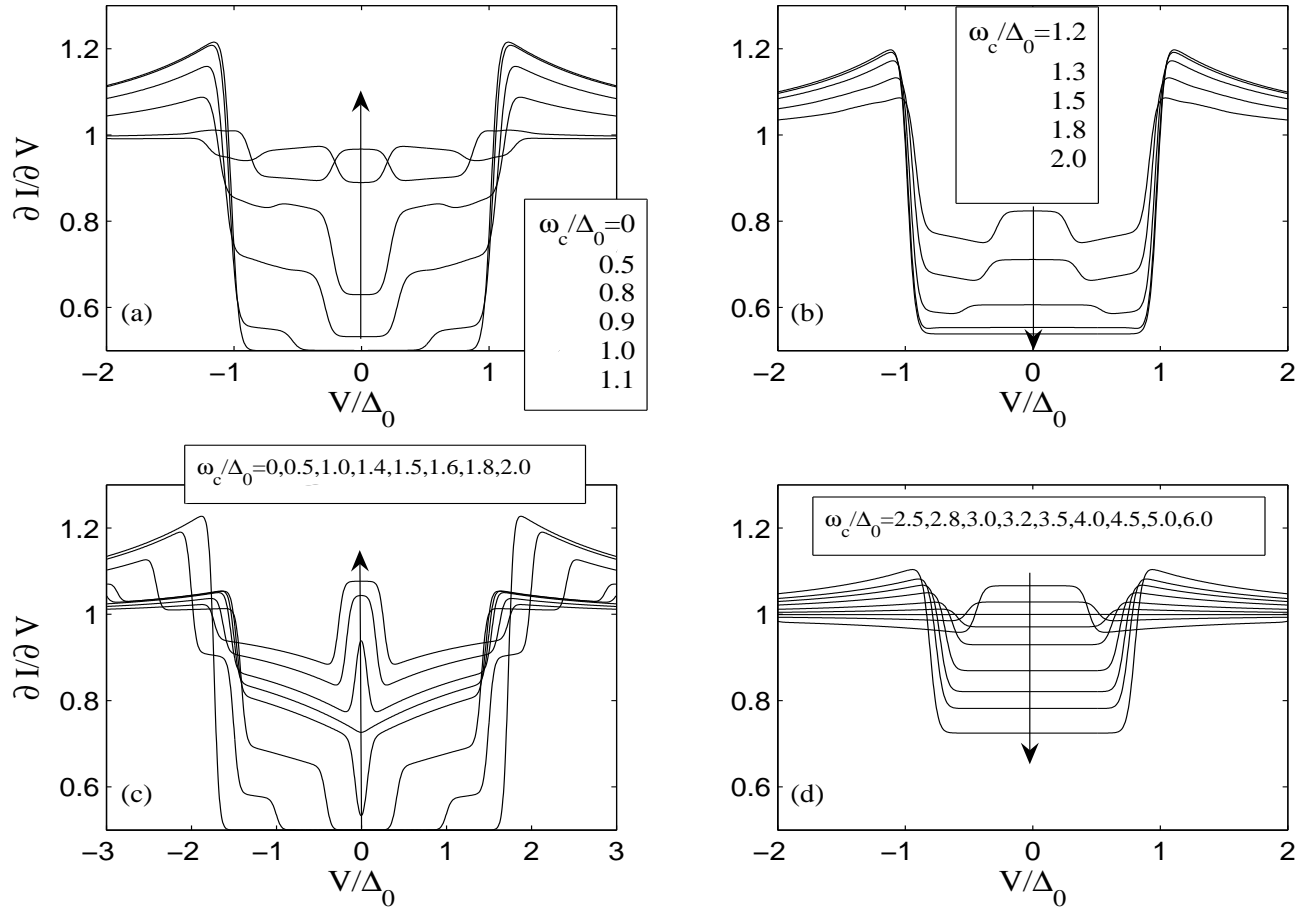


FIG. 3: The calculated differential conductance dI/dV vs. bias-voltage V/Δ_0 ($\Delta_0 = 1.0$) under weak driving CPF $\Delta/\Delta_0 = 0.2$ (a,b) and strong CPF $\Delta/\Delta_0 = \sqrt{2}$ (c,d) with several driven frequencies at a nonzero static magnetic field. The arrow indicates the direction of increase of driven frequency ω_c . The other parameters are the same as in Fig. 1.

conductance and eventually tends to the pattern in the presence of a static magnetic field alone, because the QD spin cannot, from physical point of view, catch the details of the driving field with considerably high frequency. Finally, we point out that the ZA feature is robust over a wide region of temperature (not shown here).

Conclusion.—In summary, we have presented an analytical study of the inelastic cotunneling, including the nonequilibrium spin projections and currents, in a sin-

gle QD under a static magnetic field and a perpendicular CPF, revealing a controllable I - V pattern, the transition between ZA and splitting of the differential conductance, due to *photon-assisted spin-flip* inelastic cotunneling.

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